

# CAN THE FORMATION OF X-RAY OBSCURING TORI AND JETS IN ACTIVE GALAXIES BE DETERMINED BY ONE PARAMETER?

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## ABSTRACT

A torus of reduced differential rotation can form in the inner  $\lesssim 10$ pc core of active galactic nuclei incurring a density enhancement that can account for obscuration of X-rays in Seyferts when the initial inner core to black hole mass ratio  $\gtrsim 50$ . The same density enhancement and reduction in differential rotation can also lead to dynamo growth of poloidal fields which attain a magnitude  $\sim 10^4$ G when accreted onto the central engine. As radio jet models often employ poloidal fields as agents in extracting power for the jet luminosity, we suggest that jetted AGN might require this poloidal field production. Although the poloidal field would be originally produced in the obscuring torus, jetted objects are less likely to have obscuring tori: The poloidal field would only aid in powering jet emission after it accretes with the torus matter to the central engine. Thus, only during the relatively short torus accretion time scale could there be both a jet and and torus.

**Subject Headings:** Galaxies: Active, Jets, Magnetic Fields, Radio Continuum: Galaxies

## 1. Introduction

At the centers of galaxies, and particularly in active galactic nuclei (AGN), are likely black holes (BH) which power the central luminosity by accretion (e.g., Rees 1984). Time variability and the estimated accretion efficiency seem to require black hole masses  $\gtrsim 10^7 M_\odot$  in many jetted and jet-free sources. Little is known about the detailed mass distribution in the “inner core” (IC)  $\lesssim 10$ pc regions of all galaxy types, but there in our own Galaxy the rotation curve seems to incur a change from Keplerian to flat (Genzel & Townes 1987). The transition region is quasi-rigidly rotating. This leads to a density enhancement over a purely Keplerian curve, as we show below, and may explain the presence of the circumnuclear ring (CNR) of the Galaxy and X-ray obscuring Seyfert tori (Yi et al. 1994, Duschl 1989) in accordance with unified models of AGN (e.g. Antonucci, 1993).

We show that such a region also favors the dynamo production of poloidal magnetic field (PMF) to a magnitude which, when accreted onto an AGN central engine, is likely  $\sim 10^4$ G—in agreement with that inferred by other equipartition estimates (Begelman et al. 1984). Such PMF is often required in jet models (Blandford & Znajek 1977, Lovelace et al. 1987, Appl & Camenzind 1993, Lynden-Bell 1995). The PMF dynamo growth time scale is much smaller than the torus depletion time, so whether significant PMF is produced depends only on the initial to BH mass ratio. Though the jet PMF originates in the torus, the field can only play a role in jets after accreting to the central engine. Only during the short time when the torus is depleting could there be both a jet and torus in one object. This is consistent with the fact that direct evidence for tori comes mainly from radio quiet

objects (Urry and Padovani, 1995), but more data are needed.

## 2. Estimation of Time Scales and Adiabatic Black Hole Growth

Rotation curves of spiral galaxies show quasi-rigid rotation in the inner  $\sim$  kpc, and flat rotation curves outside  $\sim$  kpc (Oort 1978, Binney & Tremaine 1987). Models which account for observed galaxy gas rotation curves seem to require (Binney & Tremaine 1987) (i) a central BH (ii) a stellar IC within a few-10pc, (iii) a more diffuse nuclear bulge of several 100pc, and (iv) an isothermal sphere of dark matter on kpc scales. Here we are interested in (i), the IC region sub-structure to the overall rotation curves, which is likely similar for all galaxy types.

We first show that a typical BH grows by accretion slowly compared to orbital time scales, but rapidly compared to the IC relaxation time, so the hole's growth is nearly adiabatic: At any time, the accreting region will be in an approximately steady state if the viscous time scale  $\tau_{vis} \ll \tau_g$ , where  $\tau_g$  is the BH growth time scale. An estimate for  $\tau_{vis}$  is

$$\tau_{vis} \sim \frac{R_d^2}{\nu} \sim \frac{R_d}{V_r} \sim \frac{R_d}{(10^{-2}V_\phi)} \sim 5 \times 10^6 \text{yr} \left( \frac{R_d}{3 \times 10^{19} \text{cm}} \right)^{3/2} \left( \frac{M_H}{10^7 M_\odot} \right)^{-1/2}, \quad (1)$$

where  $V_\phi$  is the rotational velocity,  $V_r$  is the radial velocity, and  $R_d$  is radial extent of the accretion flow. We estimate the BH growth time by assuming that the objects radiate at Eddington luminosity  $L_{Edd}$ . For a central BH of mass  $M_H$ ,  $L_{Edd} = 4\pi G(M_H/M_\odot)c/\kappa = 1.3 \times 10^{38}(M_H/M_\odot)$  erg/sec, where  $\kappa$  is the Thomson opacity  $\sim 0.4 \text{cm}^2 \text{g}^{-1}$ . As the AGN radiates at  $\sim L_{Edd}$  we have  $\dot{M}_{Edd} = 2.2 \times 10^{-9} \epsilon^{-1} M_H \text{yr}^{-1}$ , where  $\epsilon \lesssim 1$  is the efficiency

factor for radiation by accretion. Then

$$\tau_{vis} \ll \tau_g \lesssim M_H / \dot{M}_{Edd} = 4.5\epsilon \times 10^8 \text{yr}. \quad (2)$$

We must also compare  $\tau_r$ , the gravitational relaxation time of the IC region, with  $\tau_o$ , the orbital time scale. When the BH is small, its radius of influence  $R_{In} \sim GM_H/\sigma^2$  satisfies  $R_{In} < R_{IC}$ , where  $R_{IC}$  is the IC radius and  $\sigma$  is the stellar velocity dispersion, so that  $\tau_r$  is independent of  $M_H$ . When  $R_{In} \geq R_{IC}$  the relaxation time depends on  $M_H$ . In the Fokker-Planck approximation, for the case  $R_{In} < R_{IC}$  we have (Binney & Tremaine 1987)

$$\tau_r \sim 10^{11} \text{yr} \quad (\sigma/100 \text{km s}^{-1})^3 (10^5 M_\odot \text{pc}^{-3} / \rho_{IC}), \quad (3)$$

where  $\rho_{IC}$  is the IC mass density. For  $R_{In} \gg R_{IC}$  the relaxation time is

$$\begin{aligned} \tau_r &\sim \sigma^3 / (G^2 M_\odot \rho_{IC}) = \sigma^3 (4\pi R_{In}^3 / 3) / (G^2 M_\odot^2 N) = 4\pi G M_H^3 / (3\sigma^3 M_\odot^2 N) \\ &\sim 10^{13} \text{yr} \quad (M_H/10^8 M_\odot)^3 (\sigma/100 \text{km s}^{-1})^{-3} (N/10^9)^{-1}, \end{aligned} \quad (4)$$

where  $N$  is the number of IC stars.

Because of (2),(3) and (4),  $\tau_o \ll \tau_H \ll \tau_r$ , where  $\tau_H$  is the Hubble time  $\sim 10^{10}$  yr and  $\tau_o \lesssim 10^6 \text{yr} (R_{IC}/10 \text{pc}) (\sigma/100 \text{km s}^{-1})^{-1}$ , the adiabatic approach is appropriate. Young (1980) considers the adiabatic evolution of an isothermal sphere with a growing BH. Quinlan et al. (1995) confirm the results of Young (1980) and extend to non-isothermal spheres. As applied to the IC, these are reasonably consistent with a total mass dependence on radius given by

$$M_{tot}(r) = M_H [1 + m(r/R_{IC})^n], \quad r \leq R_{IC},$$

$$M_{tot}(r) = (M_H + M_{IC})r/R_{IC}, \quad r > R_{IC}, \quad (5)$$

where  $m \equiv M_{IC}/M_H$  is the ratio of BH to stellar core mass, and  $n \sim 3$ .

### 3. Gas Density

The surface gas density is given by  $\Sigma(r) = 2H(r)\rho(r)$ , where  $H(r)$  is the half-scale height of the gas and  $\rho(r)$  is the gas density. Following the standard treatment (e.g., Pringle 1981) we take the viscosity coefficient to be  $\nu = \gamma v_T H$ , where we assume  $\gamma \lesssim 1$  is the viscosity parameter and  $v_T$  is the turbulent eddy speed. Combining this with the conservation of gas mass and angular momentum, the surface density satisfies (Pringle 1981, Yi et al. 1994)

$$d\Sigma/dr + f(r)\Sigma = g(r), \quad (6)$$

where  $f(r) = (r^3\Omega'/\Omega)'/(r^3\Omega'/\Omega)$ ,  $g(r) = -[\dot{M}\Omega/(2\pi\gamma v_T^2 r)][2(\Omega/r\Omega') + 1]$ , and  $\Omega$  is the angular velocity determined by the potential. The  $'$  indicates  $d/dr$ . The solution of (6) is given by  $\Sigma(r) = \exp\left[-\int_{r_0}^r ds f(s)\right] \left(\Sigma_0 + \int_{r_0}^r ds g(s) \exp\left[\int_{r_0}^s d\lambda f(\lambda)\right]\right)$ , where  $\Sigma(r_0) = \Sigma_0$ .

To find the relationship between the surface density and the stellar mass density we note that for circular orbits,  $\Omega = (GM/r^3)^{1/2}$ . Plugging in for  $f(r)$  and  $g(r)$ , using (5),  $S \equiv r/R_{IC}$ , and  $\lambda \equiv m^{1/n}G$  we have

$$\frac{\Sigma(r)}{\Sigma_0} = \frac{(1 + mS^n)}{R_{IC}^2([n-3]mS_0^{n+2} - 3S_0^2)} \left[ \frac{R_{IC}^2([n-3]mS_0^{n+2} - 3S_0^2)}{1 + mS_0^n} - \frac{K}{\Sigma_0} \int_{m^{1/n}S_0}^{m^{1/n}S} d\lambda \frac{(n+1)\lambda^n + 1}{(\lambda + \lambda^{n+1})^{1/2}} \right], \quad (7)$$

where  $K = \dot{M}m^{\frac{1}{2}(1-1/n)}(GM_H R_{IC})^{1/2}/(2\pi\gamma v_T^2)$ .

The third term in (7) is negative, so the second term gives the correct order of magnitude. Fig. 1 shows the surface density at  $r = R_{IC} \sim R_T$ , the edge of torus, (outside of

which the potential drops) for 4 values of  $n$  as a function of  $m$  and vice versa. For large  $m$  and  $n \sim 3$  the density can be 100 times the Keplerian (i.e. when  $m \ll 1$ ) value. That  $n \sim 3$  is consistent with Young (1980) and Quinlan et al. (1995). This density enhancement can account for X-ray obscuration in Seyferts (Yi et al. 1994).

#### 4. Rigid Rotation and Poloidal Field Growth

As shown above, the enhanced density torus results from reduced differential rotation. We can calculate the reduction in  $\Omega'$  by setting  $\Omega = (GM_{tot}(r)/r^3)^{1/2}$ . Using (5) with  $n = 3$  gives  $\Omega'(R_{IC}) = -(3/2)G^{1/2}R_{IC}^{-5/2}M_H^{1/2}/(1+m)^{1/2}$ . For  $m > 50$  this gives a factor  $> 7$  reduction in  $\Omega'$  from the Keplerian value. This region may be important for the dynamo generation of PMF for radio sources as we now describe.

The mean field dynamo theory (Moffatt 1978, Parker 1979) splits the velocity and the magnetic field into mean  $(\bar{\mathbf{V}}, \bar{\mathbf{B}})$  and fluctuating  $(\mathbf{v}', \mathbf{b}')$  components. The time evolution of mean magnetic field is given by (Moffatt 1978, Parker 1979)

$$\partial \bar{\mathbf{B}} / \partial t = \nabla \times [\bar{\mathbf{V}} \times \bar{\mathbf{B}} + \alpha \bar{\mathbf{B}} - (\lambda_M + \beta) \nabla \times \bar{\mathbf{B}}], \quad (8)$$

where  $\alpha$  and  $\beta$  are the helicity and diffusion dynamo coefficients and are functions of the turbulent velocity. A non-vanishing  $\alpha$  is the result of buoyant eddies rising in an upwardly decreasing density gradient, while conserving their angular momenta. The magnetic viscosity,  $\lambda_M$ , satisfies  $\lambda_M \ll \beta$ .

For small  $\Omega'$ , the “ $\alpha^2$  dynamo” (Moffatt 1978) is favored, because the maximum growth rate depends on  $\alpha^2$  as we will see. Sufficient reduction of  $\Omega'$  means that the radial

PMF produced by the  $\alpha$  effect is not sheared into toroidal field. Simulations (e.g., Donner & Brandenburg 1990) show that dipole modes are favored in such a dynamo, in contrast to the  $\alpha - \Omega$  dynamo often applied to disks (Parker 1979).

To estimate when the  $\alpha^2$  dynamo is favored, we work in cylindrical coordinates  $(r, \phi, z)$  and write  $\bar{\mathbf{V}} = r\Omega(r, z)\hat{\mathbf{e}}_\phi$ ,  $\bar{\mathbf{B}} = \bar{B}_\phi(r, z)\hat{\mathbf{e}}_\phi + \bar{\mathbf{B}}_P$ , where  $P$  indicates the poloidal  $(r, z)$  component, and  $\phi$  indicates the toroidal component.  $\bar{\mathbf{B}}_P = \nabla \times \bar{A}(r, z)\hat{\mathbf{e}}_\phi$ . Assuming  $\alpha$  and  $\beta$  are constant, (8) can be written (Moffatt 1978)

$$\partial \bar{B}_\phi / \partial t = r(\bar{\mathbf{B}}_P \cdot \nabla)\Omega - \alpha(\nabla^2 - r^{-2})\bar{A} + \beta(\nabla^2 - r^{-2})\bar{B}_\phi, \quad (9)$$

and

$$\partial \bar{A} / \partial t = \alpha \bar{B}_\phi + \beta(\nabla^2 - r^{-2})\bar{A}. \quad (10)$$

An  $\alpha^2$  dynamo will dominate the  $\alpha - \Omega$  dynamo when the second source term on the right of (12) dominates the first, that is when  $\alpha/r \gg |r\Omega'|$ . From Parker (1979), we have  $\alpha \sim 0.4v_T$ . For a turbulently supported dust torus, observations require  $v_T/V_\phi \sim 0.5$  where  $V_\phi = r\Omega$  (Krolik & Begelman 1988, Urry & Padovani 1995). Using these and (5), the requirement near  $r = R_{IC} \sim R_T$  becomes  $0.20M_H^{1/2}(1+m)^{1/2} > (3/2)M_H^{1/2}(1+m)^{-1/2}$ , or

$$m \gg 6.5. \quad (11)$$

When (11) is satisfied, we can ignore the first term on the right of (8) near  $r = R_{IC}$ . We capture the essence of an  $\alpha^2$  dynamo, by considering the case when the  $z$ -variation dominates and assuming solutions of the form  $\bar{B}_\phi, \bar{A} \propto r \exp(\gamma t + k_z z)$ . Plugging these into (9) and (10) gives

$$\gamma \bar{B}_\phi = \alpha \bar{A} k_z^2 - \beta \bar{B}_\phi k_z^2 \quad (12)$$

and

$$\gamma\bar{A} = \alpha\bar{B}_\phi - \beta\bar{A}k_z^2, \quad (13)$$

so that the growth rate is given by

$$\text{Re}[\gamma] = -\beta k_z^2/2 + 3k_z\alpha/2. \quad (14)$$

The growth rate is positive if  $k_z < 3\alpha/\beta$ . The maximum growth rate is  $\text{Re}[\gamma]_{max} = (9/8)\alpha^2/\beta$ , showing the  $\alpha^2$  dependence. The second term on the right of (14) provides a more conservative estimate.

Let us see why an  $\alpha^2$  dynamo favors PMF. For an AGN torus, the fluctuation scale is determined by the size of dust containing clouds. The dust must be in clouds because it could not survive if the random velocities of  $\geq 100\text{km/s}$  were thermal. Thus the cloud size  $r_c$ , satisfies  $r_c < R_T$ , where the torus radius  $R_T$  is the scale for variation of the mean quantities. We can estimate the dust cloud size from observations of the CNR of our Galaxy (Genzel & Townes 1987), which is thought to be similar to the AGN dust tori (Krolik & Begelman 1988). These observations (Genzel & Townes 1987) show clouds with  $0.1 \lesssim r_c \lesssim 0.25\text{pc}$ . Now we estimate the PMF produced: Setting  $\text{Re}[\gamma] \sim k_z\alpha$  and using  $k_z\bar{A} \sim \bar{B}_P$ , with  $\alpha \sim 0.4v_T$  and  $\beta \sim (1/3)r_cv_T$  in (12) and (13) gives  $\bar{B}_P \sim \bar{B}_\phi$  for the  $\alpha^2$  dynamo. The analogous equations to (12) and (13) for the  $\alpha - \Omega$  dynamo, derived by keeping the first term on the right of (8) and dropping the term linear in  $\alpha$ , give  $[\bar{B}_P/\bar{B}_\phi]_{\alpha-\Omega} \sim r_c/R_T$ . Thus  $[B_p/\bar{B}_\phi]_{\alpha^2}/[B_p/\bar{B}_\phi]_{\alpha-\Omega} \sim R_T/r_c \gtrsim 50 - 100$ , showing that the  $\alpha^2$  dynamo favors PMF in comparison to the  $\alpha - \Omega$  dynamo. This ratio is important for jet models, particularly when the resulting luminosity depends on  $B_p^2$  (e.g., Blandford



& Znajek 1977). Thus a factor of 50-100 in the field corresponds to a factor of  $2.5 \times 10^3$  to  $10^4$  in the jet luminosity.

## 5. Poloidal Field in AGN

For PMF to be produced by a working torus dynamo, (11) must hold. In addition, the dynamo growth time must be shorter than the torus lifetime. That is,

$$\tau_d < \tau_g M_T / M_H. \quad (15)$$

Note  $\tau_d$  for the torus must be less than the field diffusion time:  $\tau_d < (10\text{pc})^2 / \beta \sim 100\text{kpc}^2 / (100\text{kms}^{-1} \cdot 10\text{kpc}) \sim 10^5\text{yr}$ . For a density  $\rho_T \sim 5 \times 10^{-18}\text{g/cm}^3$  (Krolik & Begelman 1988) and radius 5pc with height 2.5pc,  $M_T \sim 10^7 M_\odot$ . From (2) with  $\epsilon \sim 0.1$ ,  $\tau_g \sim 5 \times 10^7\text{yr}$ . Thus violating (15) requires the extreme case of  $M_H > 10^4 M_T$  so that (8) is more stringent.

The larger  $m$  is the greater the density enhancement and produced PMF. Equipartition between turbulent and magnetic energy densities gives an upper limit to the field. For  $\rho_T \sim 5 \times 10^{-18}\text{g/cm}^3$  (Krolik & Begelman 1988), corresponding to  $m \gtrsim 50$ , and  $v_T \sim 100\text{km/s}$  the turbulent energy is  $(1/2)\rho_T v_T^2 \sim 2.4 \times 10^{-4}\text{erg/cm}^3$ . Setting this equal to  $B_P^2 / 8\pi$  we have  $B_P \lesssim 8 \times 10^{-2}\text{G}$ . The field is accreted to the central engine as the torus depletes. The radial component of PMF is then sheared, but only that fraction of the field at a much lower scale height than that of the  $\sim 5 - 10\text{pc}$  torus. The z-component is unaffected by the shear. An estimate of the accreted PMF can then be made from flux freezing. For an ion-electron accretion disk with height to radius ratio  $H_d / R_d \sim 1/50$  and density  $\rho_{disk} \sim 10^{-8}\text{cm}^{-3}$  at  $r \sim 10^{14}\text{cm}$  (e.g., Celotti et al. 1992), flux conservation implies that PMF will accrete to

$B_{P,disk} < B_{P,T}[\rho_{disk}(H_d/R_d)/\rho_T(H_T/R_T)]^{2/3} \sim 8 \times 10^{-2} [10^{-8}(1/50)/5 \times 10^{-18}(1/2)]^{2/3} \sim 1.5 \times 10^4 \text{G}$ , in agreement with standard estimates (Begelman et al. 1984).

Only when the torus depletes by accretion can the field produced there move to the central engine and play its role in jets. As any jet formation time in the engine is likely much shorter than the torus depletion time, only during the latter time scale can a jetted object show both a jet and a torus. For a  $10^6 M - 10^7 M_\odot$  hole accreting at the Eddington rate with  $\epsilon \sim 0.1$  and a torus mass of  $10^7 M_\odot$ , this accretion period lasts  $\lesssim 10^7 - 10^8$  yr. Note also that a torus scale dynamo need not determine the final scale of the magnetic field, but would mediate the initial energy extraction from the rotating central source. The initial jet flow could drag the resulting field to kpc-Mpc scales as in Blandford & Rees (1974).

## 6. Conclusions

Reduced  $\Omega'$  tori in the central  $\lesssim 10$  pc regions of AGN can obscure X-rays and incur dynamo production of PMF likely required for jets, when  $m \gtrsim 50$  initially. PMF growth in the torus allows angular momentum transport, and the torus will accrete to the central engine carrying its field. PMF transport and torus depletion are associated processes and jetted objects would be less likely to show obscuring tori. Jet-free AGN would either have an initial  $m < 50$  or have their jets beamed away from us. There are a few radio loud quasars (RLQ) or Seyferts with inferred BH masses  $> 10^8 M_\odot$  (e.g. NGC 5548, Krolik et al. 1991). The absence of direct evidence for dust tori in the latter may be consistent with our paradigm, but may require non-adiabatic analysis.

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### Figure Caption

Figure 1: (a) Surface density at the torus' outer edge as function of  $m$  for  $n = 0.01, 1, 2.7, 3$  going from the bottom to the top curves. (b) Surface density as a function of  $n$  for  $m = 0.01, 1, 10, 100$  from the bottom to top curves.